

## Bulk U(1) Messenger

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### Abstract

We propose a new U(1) gauge interaction in the bulk in higher dimensional space-time, which transmits supersymmetry-breaking effects on the hidden brane to the observable our brane. We find that rather small gauge coupling constant of  $U(1)_{\text{bulk}}$ ,  $\alpha_{\text{bulk}} \simeq 5 \times 10^{-4}$ , is required for a successful phenomenology. This result implies the compactification length  $L$  of the extra dimension to be  $L^{-1} \simeq 2 \times 10^{15}$  GeV for (4+1)-dimensional spacetime. This large compactification length  $L$  is a crucial ingredient to suppress unwanted flavor-changing neutral currents and hence our proposal is very consistent with the Randall-Sundrum brane-world scenario.

Most of supersymmetry (SUSY) breaking models in supergravity (even including gauge mediation models) assume a separation of the SUSY-breaking and the SUSY standard-model sectors [1, 2]. However, the origin of the separation is not well understood, although such a separation is crucial to obtain phenomenologically consistent spectra for SUSY particles.

The brane world proposed by Randall and Sundrum [3] provides a beautiful geometric explanation for the separation. That is, the hidden and observable sectors live on different three-dimensional branes separated by a gravitational bulk [4] in higher dimensional space-time. It has been, recently, claimed [5, 3] that this brane separation produces the hidden and observable separation in the “conformal” frame in supergravity, which was proposed long time ago from a phenomenological ground [6].

It is a crucial observation in Ref. [6] that the above separation in the “conformal” frame induces a no-scale type Kähler potential<sup>1</sup> in the Einstein frame,

$$K(\Phi_{\text{obs}}, \Phi_{\text{obs}}^\dagger, \Phi_{\text{hid}}, \Phi_{\text{hid}}^\dagger) = -3 \log \left( 1 - \frac{1}{3} f_O(\Phi_{\text{obs}}, \Phi_{\text{obs}}^\dagger) - \frac{1}{3} f_H(\Phi_{\text{hid}}, \Phi_{\text{hid}}^\dagger) \right). \quad (1)$$

Here,  $\Phi_{\text{obs}}$  and  $\Phi_{\text{hid}}$  denote superfields in the observable and hidden sectors, respectively. With the above Kähler potential Eq. (1) we easily show [6] that all soft SUSY-breaking masses and  $A$  terms in the observable sector vanish in the limit of the zero cosmological constant. All gaugino masses in the observable sector also vanish because of the decoupling of the hidden superfield  $Z$  from the gauge kinetic function [3].<sup>2</sup>

On the contrary, the SUSY-invariant  $\mu$  term ( $\mu H \bar{H}$ ) naturally arises [9] from the Kähler potential if  $f_O$  contains  $f_O \supset H \bar{H}$ , where  $H$  and  $\bar{H}$  are chiral superfields of Higgs doublets. This mechanism [9] produces the SUSY-invariant mass  $\mu$  of the order of the gravitino mass, *i.e.*  $\mu \simeq m_{3/2}$ . This requires the gravitino mass  $m_{3/2} \simeq 100 \text{ GeV} - 1 \text{ TeV}$  for the correct electroweak symmetry breaking. With these gravitino masses the anomaly mediation [3, 10]

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<sup>1</sup> The no-scale supergravity [7] adopts a specific form  $f_H = Z + Z^\dagger$ , where  $Z$  is the superfield responsible for the SUSY breaking. We assume the Kähler potential for the  $Z$  field to be of the form  $f_H = Z^\dagger Z + \dots$ , where the ellipsis denotes higher order terms.

<sup>2</sup> The gaugino-mediated SUSY breaking was proposed in Ref. [6]. See also recent works [8].

generates too small SUSY-breaking masses in the observable sector and hence all SUSY particles in the observable sector except for the Higgsinos and the gravitino remain almost massless.

In this letter we introduce a new  $U(1)$  gauge interaction in the bulk to solve the above problem.<sup>3</sup> The  $U(1)_{\text{bulk}}$  gauge superfield plays a role of messenger between the hidden and the observable branes. SUSY-breaking effects on the hidden brane are transmitted to the observable brane through the bulk  $U(1)_{\text{bulk}}$  gauge interaction and all of the gauginos, squarks and sleptons in the observable sector acquire suitable SUSY-breaking masses.

It is well known that a similar  $U(1)$  gauge interaction is also used as a messenger between the SUSY-breaking and observable sectors in a class of gauge-mediated SUSY breaking models [12, 13]. Thus, it is quite natural to identify the above bulk  $U(1)_{\text{bulk}}$  gauge interaction with the messenger  $U(1)_m$  gauge interaction in the gauge mediation models. We adopt a model proposed in Ref. [13], and interpret it as a low-energy effective theory of the brane world. We then show that rather small gauge coupling  $\alpha_{\text{bulk}} \simeq 5 \times 10^{-4}$  of the  $U(1)_{\text{bulk}}$  is required for a successful phenomenology. This small coupling is regarded as a consequence of a large volume of extra dimension. In fact, the result implies the compactification length  $L$  of the extra dimension to be  $L^{-1} \simeq 2 \times 10^{15}$  GeV for  $(4 + 1)$ -dimensional spacetime and the fundamental scale  $M_*$  is determined as  $M_* \simeq 2 \times 10^{17}$  GeV to reproduce the gravitational scale  $M_G \simeq (2M_*^3 L)^{1/2} \simeq 2 \times 10^{18}$  GeV. Such a large compactification length is a crucial ingredient to suppress the flavor-changing neutral currents (FCNC's) and hence our proposal is very consistent with the brane-world scenario of Randall and Sundrum [3].

Let us first briefly review the gauge mediation model proposed in Ref. [13]. The model consists of three sectors: dynamical SUSY breaking (DSB) sector, messenger sector, and the minimal SUSY standard model (MSSM) sector. The DSB sector is based on a SUSY  $SU(2)$  gauge theory with four doublet chiral superfields  $Q_i$  where  $i$  is a flavor index ( $i = 1, \dots, 4$ ). Here, we have a global flavor  $SU(4)_F$ . We assume the following tree-level superpotential

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<sup>3</sup> A similar model in the brane world has been also discussed in Ref. [11].

introducing six singlet chiral superfields  $Z$  and  $Z^a$  ( $a = 1, \dots, 5$ ):

$$W_{\text{tree}} = \lambda Z(QQ) + \lambda_Z Z^a (QQ)_a, \quad (2)$$

where  $(QQ)$  and  $Z$  are singlets of the  $\text{SP}(4)_F$  subgroup of the flavor  $\text{SU}(4)_F$  and  $(QQ)_a$  and  $Z^a$  are five-dimensional representations of the  $\text{SP}(4)_F$ . As shown in Ref. [14], integration of the  $\text{SU}(2)$  gauge fields together with  $Q_i$  and  $Z^a$  leads to the low-energy effective superpotential

$$W_{\text{eff}} \simeq \frac{\lambda}{(4\pi)^2} \Lambda^2 Z, \quad (3)$$

for  $\lambda_Z > \lambda$ , where  $\Lambda$  is a dynamical scale of the  $\text{SU}(2)$  gauge interaction.<sup>4</sup> We have nonvanishing  $F$  term,  $\langle F_Z \rangle \simeq \lambda \Lambda^2 / (4\pi)^2 \neq 0$ , and hence SUSY is broken [16]. We also assume that the fields in the DSB sector are charged under the  $\text{U}(1)_m$  gauge interaction. The charge assignments of chiral superfields are given by [13]

$$Q_1(+1), \quad Q_2(-1), \quad Q_3(0), \quad Q_4(0). \quad (4)$$

Here, the numbers in each parentheses denote the  $\text{U}(1)_m$  charges. The  $\text{U}(1)_m$  charges for  $Z$  and  $Z^a$  are determined such that the superpotential Eq. (2) is invariant under the  $\text{U}(1)_m$ .

We now turn to the messenger sector. It consists of three chiral superfields,

$$E(+1), \quad \bar{E}(-1), \quad S(0), \quad (5)$$

and vector-like messenger quark and lepton superfields,  $d, \bar{d}, l, \bar{l}$ . Here, the messenger quark multiplets  $d, \bar{d}$  and lepton multiplets  $l, \bar{l}$  are all neutral under the  $\text{U}(1)_m$ . The  $d$  and  $\bar{d}$  ( $l$  and  $\bar{l}$ ) transform as the right-handed down quark and its antiparticle (the left-handed lepton doublet and its antiparticle) under the standard-model gauge group, respectively. The superpotential for the messenger sector is given by

$$W_{\text{mess}} = k_E S E \bar{E} + \frac{f}{3} S^3 + k_d S d \bar{d} + k_l S l \bar{l}. \quad (6)$$

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<sup>4</sup> The factor of  $4\pi$  is determined by the naïve dimensional analysis [15].

The SUSY-breaking effects in the DSB sector are transmitted to the messenger sector through the  $U(1)_m$  gauge interaction. As a result, the  $E$  and  $\bar{E}$  fields obtain positive soft SUSY-breaking squared masses  $m_E^2$  and  $m_{\bar{E}}^2$ ,<sup>5</sup>

$$m_E = m_{\bar{E}} \sim \frac{\alpha_m}{4\pi} \frac{\lambda F_Z}{\Lambda} \simeq \frac{\alpha_m}{4\pi} \frac{\lambda^2}{16\pi^2} \Lambda, \quad (7)$$

where  $\alpha_m = g_m^2/4\pi$  is the  $U(1)_m$  gauge coupling constant. They generate the following negative soft SUSY-breaking mass squared  $-m_S^2$  for the  $S$  field through the Yukawa coupling  $k_E S E \bar{E}$  in Eq. (6) at the one-loop level:

$$m_S^2 \simeq \frac{4}{(4\pi)^2} k_E^2 m_E^2 \ln \frac{\Lambda}{m_E}. \quad (8)$$

Altogether, the resulting scalar potential is

$$V_{\text{mess}} = \sum_{\eta} \left| \frac{\partial W_{\text{mess}}}{\partial \eta} \right|^2 + m_E^2 |E|^2 + m_{\bar{E}}^2 |\bar{E}|^2 - m_S^2 |S|^2, \quad (9)$$

where  $\eta$  denotes chiral superfields  $E, \bar{E}, S, d, \bar{d}, l$  and  $\bar{l}$ . This potential has a global minimum at

$$\begin{aligned} \langle S^* S \rangle &= \frac{m_S^2}{2f^2}, & \langle |F_S| \rangle &= \frac{m_S^2}{2f}, \\ \langle E \rangle &= \langle \bar{E} \rangle = \langle d \rangle = \langle \bar{d} \rangle = \langle l \rangle = \langle \bar{l} \rangle = 0, \end{aligned} \quad (10)$$

in a certain parameter region [13]. Thus, all the standard-model gauge symmetries are preserved and the SUSY-breaking effects are transmitted to the messenger quark and lepton multiplets through  $\langle F_S \rangle$ .

The soft SUSY-breaking masses for the gauginos  $\tilde{g}_i$  ( $i = 1, \dots, 3$ ) and the squarks, sleptons, and Higgses  $\tilde{f}$  in the MSSM sector are generated by integrating out the messenger quarks and leptons as

$$m_{\tilde{g}_i} = c_i \frac{\alpha_i}{4\pi} \Lambda_{\text{mess}}, \quad (11)$$

$$m_{\tilde{f}}^2 = 2\Lambda_{\text{mess}}^2 \left[ C_3 \left( \frac{\alpha_3}{4\pi} \right)^2 + C_2 \left( \frac{\alpha_2}{4\pi} \right)^2 + \frac{5}{3} Y^2 \left( \frac{\alpha_1}{4\pi} \right)^2 \right], \quad (12)$$

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<sup>5</sup> If  $\langle Z \rangle = 0$ , there is an unbroken  $U(1)$   $R$ -symmetry. In this case, the  $U(1)_m$  gaugino remains massless while  $E$  and  $\bar{E}$  fields obtain soft SUSY-breaking masses in Eq. (7).

where  $c_1 = 5/3$ ,  $c_2 = c_3 = 1$ ;  $C_3 = 4/3$  for color triplets and zero for singlets,  $C_2 = 3/4$  for weak doublets and zero for singlets, and  $Y$  is the hypercharge ( $Y = Q_{\text{em}} - T_3$ ). Here,  $\Lambda_{\text{mess}}$  is an effective messenger scale defined as

$$\Lambda_{\text{mess}} \equiv \frac{\langle |F_S| \rangle}{\langle |S| \rangle} = \frac{m_S}{\sqrt{2}}, \quad (13)$$

which can be written in terms of the SUSY-breaking scale  $\sqrt{F_Z}$  as

$$\begin{aligned} \Lambda_{\text{mess}} &\simeq \frac{\sqrt{2}}{(4\pi)^4} \alpha_m \lambda^2 k_E \sqrt{\ln \frac{(4\pi)^3}{\alpha_m \lambda^2}} \cdot \Lambda \\ &= \frac{\sqrt{2}}{(4\pi)^3} \alpha_m \lambda \sqrt{\lambda} k_E \sqrt{\ln \frac{(4\pi)^3}{\alpha_m \lambda^2}} \sqrt{F_Z}. \end{aligned} \quad (14)$$

We are now at the point of this letter. We consider that the above model is the low-energy effective theory of the brane world, in which all fields in the DSB sector reside on the hidden brane while the messenger and MSSM sectors are localized on the observable brane. Then, the  $U(1)_m$  gauge multiplet should necessarily live in the bulk. Thus, we identify  $U(1)_m$  with the bulk  $U(1)_{\text{bulk}}$ .<sup>6</sup> The SUSY breaking on the hidden brane is transmitted to the observable brane by the  $U(1)_{\text{bulk}}$  gauge and gravitational interactions across the bulk between two branes.

The effective messenger scale  $\Lambda_{\text{mess}}$  should be taken at  $(10^4 - 10^5)$  GeV to induce the MSSM gaugino and sfermion masses of the electroweak scale. On the other hand, the supergravity effects generate simultaneously the  $\mu$  term and the gravitino mass as discussed in the introduction [9],

$$\mu \simeq m_{3/2} = \frac{F_Z}{\sqrt{3}M_G}. \quad (15)$$

We should set  $\sqrt{F_Z} \simeq (2 - 6) \times 10^{10}$  GeV to reproduce correctly the electroweak symmetry breaking (*i.e.*  $\mu \simeq 100$  GeV – 1 TeV). It is a crucial point that the above two conditions determine the  $U(1)_{\text{bulk}}$  gauge coupling constant through Eq. (14). Assuming the Yukawa coupling constants  $\lambda$  and  $k_E$  connecting fields on the same brane to be of order one, we

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<sup>6</sup> We assume that the compactification scale  $L^{-1}$  is sufficiently high as  $L^{-1} \gg \Lambda$ .

obtain the  $U(1)_{\text{bulk}}$  gauge coupling  $\alpha_{\text{bulk}} \simeq 5 \times 10^{-4}$  at the scale  $\Lambda$ . Since the running effect of the  $U(1)_{\text{bulk}}$  gauge coupling is negligible for the matter content in the present model, we find that  $\alpha_{\text{bulk}} \simeq 5 \times 10^{-4}$  at the compactification scale of the extra dimension.

For a definiteness, we here assume the  $(4 + 1)$ -dimensional spacetime with one extra dimension compactified on the orbifold  $S^1/\mathbf{Z}_2$ . The hidden and observable branes are located at two different fixed points of the  $S^1/\mathbf{Z}_2$  separated by a distance  $L$ . In this case, the 5-dimensional  $U(1)_{\text{bulk}}$  multiplet is composed of a vector field, a Dirac spinor field and a real scalar field, which corresponds to a  $\mathcal{N} = 2$  vector multiplet in 4 dimensions. Through an orbifold projection, however, only the 4-dimensional  $\mathcal{N} = 1$  vector multiplet of the  $U(1)_{\text{bulk}}$  can couple to the fields on two branes.<sup>7</sup>

The 4-dimensional gauge coupling  $g_{\text{bulk}}$  is obtained from the 5-dimensional coupling  $g_{\text{bulk}}^{(5)}$  as

$$\frac{1}{g_{\text{bulk}}^2} = \frac{2L}{(g_{\text{bulk}}^{(5)})^2}. \quad (16)$$

Assuming that  $g_{\text{bulk}}^{(5)}$  is of order one in the unit of the fundamental scale  $M_*$ , we obtain the following relation:

$$(2L)^{-1} = (4\pi\alpha_{\text{bulk}})M_*, \quad (17)$$

which shows that the compactification length  $L$  is  $(8\pi\alpha_{\text{bulk}})^{-1} \simeq 100$  times larger than the fundamental length scale  $M_*^{-1}$  of the theory. This sufficiently suppresses the unwanted FCNC's which would be induced by exchanges of bulk fields of masses around  $M_*$  [3].

On the other hand, the gravitational scale  $M_G$  is given by

$$M_G^2 = 2M_*^3 L. \quad (18)$$

Thus, together with Eq. (17), we find that the fundamental scale  $M_*$  and the compactification scale  $L^{-1}$  are given by

$$M_* = M_G(4\pi\alpha_{\text{bulk}})^{1/2} \simeq 2 \times 10^{17} \text{ GeV}, \quad (19)$$

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<sup>7</sup> A detailed analysis on the orbifold  $S^1/\mathbf{Z}_2$  is given in Ref. [17].

$$L^{-1} = 2M_G(4\pi\alpha_{\text{bulk}})^{3/2} \simeq 2 \times 10^{15} \text{ GeV}. \quad (20)$$

It is very interesting that these values are close to the ones discussed in Ref. [4].

Several comments are in order. First, the  $U(1)_{\text{bulk}}$  gauge symmetry is unbroken in the present model so that one of scalar components of the  $E$  and  $\bar{E}$  fields is completely stable. This requires the reheating temperature  $T_R$  of inflation to be lower than the mass of the lightest scalar field of order  $\langle S \rangle \simeq 10^5 \text{ GeV}$ , in order for its present energy density not to exceed the critical density of the universe. Second, the dangerous  $D$ -term for the  $U(1)_{\text{bulk}}$  does not appear in the model [13], since there is an unbroken charge conjugation symmetry defined as

$$\begin{aligned} Q_1 &\rightarrow Q_2, & Q_2 &\rightarrow -Q_1, \\ V_{\text{bulk}} &\rightarrow -V_{\text{bulk}}, & E &\rightarrow \bar{E}, & \bar{E} &\rightarrow E, \end{aligned} \quad (21)$$

where  $V_{\text{bulk}}$  is the  $U(1)_{\text{bulk}}$  gauge superfield. The singlets  $Z$  and  $Z^a$  are assumed to transform properly so that the superpotential Eq. (2) is invariant under the charge conjugation symmetry Eq. (21). This is a consequence of the vector-like structure of the  $U(1)_{\text{bulk}}$  gauge sector. Third, if there exists a non-Abelian gauge theory in the bulk other than the  $U(1)_{\text{bulk}}$ , the radius of the extra dimension is stabilized as shown in Ref. [5].<sup>8</sup> Note that the stabilization is not disrupted by the Casimir energy induced by the SUSY breaking [17].

Finally, we should stress that the gravitino has a mass of order  $100 \text{ GeV} - 1 \text{ TeV}$  in the present model although the mass spectrum of the other SUSY particles is the same as that of the gauge-mediated SUSY breaking models.<sup>9</sup> This leads to an interesting phenomenological consequence. That is, the bino is most likely the dark matter in the present universe while the usual gauge mediation models predict the gravitino dark matter. It is a generic feature [11] of the gauge-mediated SUSY breaking models in which the phenomenologically viable  $\mu$  term arises from the supergravity effect [9]. We should stress that the Randall-Sundrum

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<sup>8</sup> The radius can also be stabilized by the mechanism of Ref. [18], if there are two non-Abelian gauge theories in the bulk which couple to suitable matters on a brane.

<sup>9</sup> The  $B$  term is of the order of the gravitino mass, where  $B$  is defined as  $\mathcal{L} = \mu B \tilde{H} \tilde{\bar{H}} + \text{h.c.}$  ( $\tilde{H}$  and  $\tilde{\bar{H}}$  are the scalar components of  $H$  and  $\bar{H}$ ). Thus, the small  $\tan\beta$  region may be accommodated in the present model.



brane-world scenario [3] provides a natural solution to the  $\mu$  problem in a large class of gauge-mediated SUSY breaking models [12, 13], although we have adopted a specific model in Ref. [13] to demonstrate our point.

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## References

- [1] For a review, H.P. Nilles, *Phys. Rept.* **110** (1984) 1.
- [2] For a review, G.F. Giudice and R. Rattazzi, *Phys. Rept.* **322** (1999) 419.
- [3] L. Randall and R. Sundrum, *Nucl. Phys.* **B557** (1999) 79.
- [4] P. Horava and E. Witten, *Nucl. Phys.* **B460** (1996) 506; *Nucl. Phys.* **B475** (1996) 94;  
E. Witten, *Nucl. Phys.* **B471** (1996) 135.
- [5] M.A. Luty and R. Sundrum, hep-th/9910202.
- [6] K. Inoue, M. Kawasaki, M. Yamaguchi and T. Yanagida, *Phys. Rev.* **D45** (1992) 328.
- [7] For a review, A. Lahanas and D.V. Nanopoulos, *Phys. Rept.* **145** (1987) 1.
- [8] D.E. Kaplan, G.D. Kribs and M. Schmaltz, hep-ph/9911293;  
Z. Chacko, M.A. Luty, A.E. Nelson and E. Ponton, *JHEP* **0001** (2000) 003.
- [9] G.F. Giudice and A. Masiero, *Phys. Lett.* **B206** (1988) 480.
- [10] G.F. Giudice, M.A. Luty, H. Murayama and R. Rattazzi, *JHEP* **9812** (1998) 027.
- [11] Izawa K.-I., Y. Nomura and T. Yanagida, *Prog. Theor. Phys.* **102** (1999) 1181.
- [12] M. Dine, A.E. Nelson and Y. Shirman, *Phys. Rev.* **D51** (1995) 1362;  
M. Dine, A.E. Nelson, Y. Nir and Y. Shirman, *Phys. Rev.* **D53** (1996) 2658.
- [13] Y. Nomura, K. Tobe and T. Yanagida, *Phys. Lett.* **B425** (1998) 107.
- [14] Izawa K.-I., Y. Nomura, K. Tobe and T. Yanagida, *Phys. Rev.* **D56** (1997) 2886.
- [15] M.A. Luty, *Phys. Rev.* **D57** (1998) 1531;  
A.G. Cohen, D.B. Kaplan and A.E. Nelson, *Phys. Lett.* **B412** (1997) 301.
- [16] Izawa K.-I. and T. Yanagida, *Prog. Theor. Phys.* **95** (1996) 829;  
K. Intriligator and S. Thomas, *Nucl. Phys.* **B473** (1996) 121.
- [17] E.A. Mirabelli and M.E. Peskin, *Phys. Rev.* **D58** (1998) 065002.
- [18] Izawa K.-I. and T. Yanagida, *Prog. Theor. Phys.* **101** (1999) 171.